

# Defect-unbinding and the Bose-glass transition in layered superconductors

C.J. van der Beek<sup>1</sup>, M. Konczykowski<sup>1</sup>, A.V. Samoilov<sup>2</sup>, N. Chikumoto<sup>3</sup>, S. Bouffard<sup>4</sup>, and M.V. Feigel'man<sup>5</sup>

<sup>1</sup>*Laboratoire des Solides Irradiés, Ecole Polytechnique, 91128 Palaiseau, France*

<sup>2</sup>*Condensed Matter Physics 114-36, California Institute of Technology, Pasadena CA 91125, U.S.A.*

<sup>3</sup>*Superconductivity Research Laboratory, ISTEK, Minato-ku, Tokyo 105, Japan*

<sup>4</sup>*Centre Interdisciplinaire de Recherche Ions Lasers (C.I.R.I.L.), B.P. 5133, 14040 Caen Cedex, France*

<sup>5</sup>*Landau Institute of Theoretical Physics, Moscow, Russia*

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The low-field Bose-glass transition temperature in heavy-ion irradiated  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$  increases progressively with increasing density of irradiation-induced columnar defects, but saturates for densities in excess of  $1.5 \times 10^9 \text{ cm}^{-2}$ . The maximum Bose-glass temperature corresponds to that above which diffusion of two-dimensional pancake vortices between different vortex lines becomes possible, and above which the “line-like” character of vortices is lost. We develop a description of the Bose-glass line that is in excellent quantitative agreement with the experimental line obtained for widely different values of track density and material parameters.

Heavy-ion irradiated (HII) layered superconductors [1,2] have recently been at the focus of attention, because the irradiation-induced amorphous columnar tracks help overcome the detrimental effects of the high material anisotropy [3–7], and partially re-establish long-range superconducting phase order [8]. In the layered superconductor  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ , heavy-ion irradiation increases the irreversibility field  $B_{irr}(T)$  below which the  $I(V)$ -curve is nonlinear due to vortex pinning on the tracks [9] to values well above the field  $B_{FOT}(T)$  at which the first order vortex lattice-to liquid transition field takes place in the pristine material [10]. At inductions  $B_{FOT}(T) \ll B \ll B_{irr}(T)$ , the irradiated superconductor displays the phenomenology of the Bose-glass phase of localized vortices [4,9,11]; moreover, the transport properties show a distinct anisotropy related to the presence of the tracks [3–6] that is absent in the pristine material [7], and which suggests that the vortices behave as well-defined separate *lines*, *i.e.* vortex lines in the Bose-glass phase are *disentangled*.

The position of  $B_{irr}(T)$  and the occurrence of flux-line entanglement were shown to be intimately related in moderately anisotropic HII superconductors such as  $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$  [12]. There,  $B_{irr}(T) \sim B_{BG}(T)$  corresponds to the second order phase transition line between the Bose-glass and the vortex liquid [11–13];  $B_{BG}(T)$  progressively increases with increasing columnar defect density  $n_d$ , to an upper limit attained when  $n_d \approx 1 \times 10^{11} \text{ cm}^{-2}$  (corresponding to an ion dose-equivalent “matching” field  $B_\phi \equiv \Phi_0 n_d = 2 \text{ T}$ ). Departing from the correspondence of vortex lines with the world lines of interacting bosons in two dimensions (2D) [11,14], it was argued [12] that the upper limit of  $B_{BG}(T, n_d)$  corresponds to the field beyond which the (entangled) vortex liquid becomes stable with respect to the introduction of linear defects. In this temperature and field regime, these produce a rather *weak* pinning due to intervortex repulsion [15,16] and the averaging effect of vortex ther-

mal excursions [14,17]. The situation in layered superconductors such as  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$  is *a priori* different. First, the weak coupling between adjacent  $\text{CuO}_2$  bilayers (with separation  $s \approx 1.5 \text{ nm}$ ) implies that vortex lines are extremely soft. Over the larger part of the first Brillouin zone of the vortex lattice, the contribution of the dipole interaction between “pancake” vortices in adjacent bilayers to the tilt modulus  $c_{44}$  is expected to exceed that of the line tension  $\varepsilon_1$ , which is determined by the interlayer Josephson effect [18,19]. Flux lines are then better described as stacks of pancakes and the analogy with the 2D boson system is no longer valid [20]. Another consequence is that pinning near  $B_{irr}(T)$  in HII  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$  is not weak: at fields  $B \lesssim \frac{1}{6} B_\phi$  pancake vortices gain maximum free energy by remaining localized on the columnar defects, even in the vortex liquid phase [21]. Nevertheless, a well-defined transition from very slow nonlinear vortex dynamics in the Bose-glass [4,9], to Ohmic response in the vortex liquid [5,6] does exist near the irreversibility line in  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ .

In order to cast light on the mechanism of this transition, we have measured  $B_{irr}(T)$  in HII  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$  for widely varying track density and oxygen content. The latter determines the values of the anisotropy parameter  $\gamma$  and the penetration depth  $\lambda_{ab}(T)$  [22]: both increase as one decreases the oxygen content towards optimal doping. It turns out that the phenomenological behavior of  $B_{irr}(T)$  at low fields is rather similar to that found in  $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ .  $B_{irr}(T)$  increases progressively towards higher values with increasing defect density, but saturates for  $n_d \gtrsim 1.5 \times 10^9 \text{ cm}^{-2}$ , or  $B_\phi \gtrsim 30 \text{ mT}$ . For higher matching fields, the low-field portion ( $B \lesssim \frac{1}{6} B_\phi$ ) of  $B_{irr}(T)$  adopts a strictly exponential temperature dependence; although  $T_{irr}(B)/T_c$  strongly increases when the value of  $\lambda_{ab}$  decreases, we will see that this increase is not due to the increase of the pinning energy.

$\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$  crystals were grown at the University of Tokyo using the travelling-solvent floating-zone

method, and then postannealed at either 800°C or 500°C. This produces  $T_c$ 's of 89 K (optimal doping) and 82 K (overdoped) respectively [22]. The crystals were cut to rectangles of size 500 (l)  $\times$  400 (w)  $\times$  20 (t)  $\mu\text{m}^3$ , and irradiated with varying fluences of 5.8 GeV Pb ions at the Grand Accélérateur National d'Ions Lourds (GANIL) at Caen, France. The ion beam was directed parallel to the sample  $c$ -axis. Each ion impact created an amorphous columnar track of radius 3.5 nm traversing the sample along its entire thickness. Samples were prepared with track densities  $1 \times 10^9 \text{ cm}^{-2} < n_d < 2 \times 10^{11} \text{ cm}^{-2}$ , corresponding to  $20 \text{ mT} \leq B_\phi \leq 4 \text{ T}$ . The irradiation caused  $T_c$  to decrease:  $\partial T_c / \partial n_d = -3.6 \times 10^{11} \text{ K cm}^2$  or  $\partial T_c / \partial B_\phi = -1.8 \text{ K T}^{-1}$ .

Subsequent measurements were performed using the Local Hall Probe Magnetometer in AC mode [4,23]. A small ac field of amplitude  $h_{ac} = 1 \text{ G}$  and frequency  $f = 7.75 \text{ Hz}$  is applied parallel to the sample  $c$ -axis, colinearly with the DC field used to create the vortices. The ac field leads to a periodic electric field gradient of magnitude  $\sim 2\pi\mu_0 h_{ac} f$  across the sample. Using a miniature Hall probe one measures the RMS induction  $B_{ac}(f, T)$  at the center of the sample top surface, which is simply related to the sample screening current [23,25].

Figure 1 shows the fundamental and third harmonic transmittivities, defined as  $T_H \equiv [B_{ac}(f, T) - B_{ac}(f, T \ll T_c)]/B$  and  $T_{H3} \equiv B_{ac}(3f, T)/B$  respectively, with

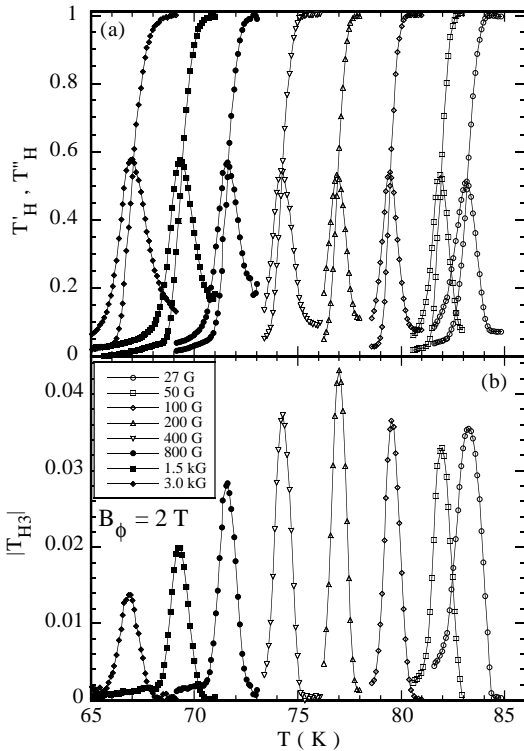


FIG. 1. Fundamental (a) and third harmonic (b) transmittivity of the  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$  single crystal with matching field  $B_\phi = 2 \text{ T}$ , for DC fields (applied parallel to the  $c$ -axis and to the columnar defects) between 27 G and 3.0 kG.

$B \equiv B_{ac}(f, T \gg T_c) - B_{ac}(f, T \ll T_c)$  [23], measured for the optimally doped crystal with  $B_\phi = 2 \text{ T}$ . The presence of a third harmonic response  $|T_{H3}|$  implies the non-linearity of the sample's  $I(V)$ -characteristic. The fields  $B_{irr}(T)$ , or temperatures  $T_{irr}(B)$ , below which the third harmonic signal can be first observed upon cooling are plotted in Fig. 2. At  $B_{irr}$ , the working point enters the nonlinear part of the sample  $I(V)$  curve at an electric field of the order of  $10^{-7} \text{ V m}^{-1}$ ; this corresponds to a voltage drop of  $\sim 50 \text{ pV}$  across the sample. At such low voltages, the measurement of the ac screening current may be sensitivity-limited [4,25]. In practice, the use of  $h_{ac} = 1 \text{ Oe}$  means that the minimum measurable current density  $j_{min} \approx 5 \times 10^2 \text{ A m}^{-2}$ , comparable to a transport current of 0.1 mA. The coincidence of our  $B_{irr}(T)$  with the  $B_{BG}$ -data determined by Seow *et al.* [5] indicates that for all practical purposes  $B_{irr}$  is a good approximation of the Bose-glass transition field (Fig. 2).

The evolution of  $B_{irr}(T)$  with  $B_\phi$  is plotted in Fig. 3. For  $B_\phi \lesssim B_\phi^{min} = 30 \text{ mT}$ ,  $B_{irr}(T)$  increases monotonically with increasing  $B_\phi$  [24]. As long as  $B_{irr}(T) < B_\phi$ , it depends exponentially on temperature, while for  $B_{irr} > B_\phi$ ,  $B_{irr} \propto (1 - T/T_c)$  [25]. For large  $B_\phi \gtrsim B_\phi^{min}$ , one can distinguish *three* distinct sections of  $B_{irr}(T)$ . At all but the very lowest fields,  $B_{irr}$  again depends exponentially on  $T$ , but with the specificity that it is *independent* of  $B_\phi$ . In this regime [(I) in Fig. 2], there thus exists an *upper limit* of  $B_{BG}(T)$  in  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$ , as is the case in  $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$  [12]. At  $B_{irr} \equiv B_{int} \sim \frac{1}{6}B_\phi$ , the exponential increase abruptly changes into a nearly vertical rise (regime II);  $B_{int}$  is the field at which intervortex repulsion start to determine the pinned vortex configuration [21]. This transition is also manifest in the vortex liquid phase as that at which a “recoupling” transition was measured using Josephson Plasma Resonance (JPR) [8]. As  $B_{irr}$  increases to a sizeable (but not constant) fraction of  $B_\phi$ , a weak temperature dependence is once

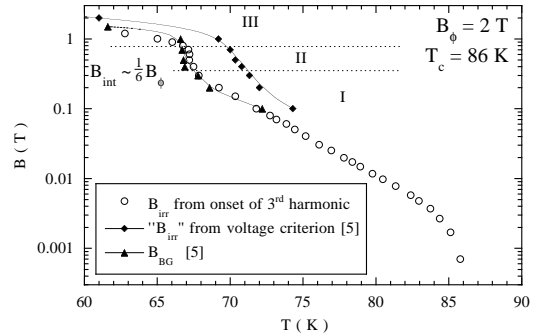


FIG. 2. The field  $B_{irr}(T)$  below which a third harmonic transmittivity signal can be observed in an optimally doped HII  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$  single crystal ( $B_\phi = 2 \text{ T}$ ) (○). Filled triangles denote the Bose-glass transition line determined from the power-law scaling of the resistivity,  $\rho \propto (T - T_{BG})^s$  of a crystal with the same  $T_c$  and  $B_\phi$  [5]; filled diamonds denote the “irreversibility line” determined from the onset of a measurable resistivity [5].

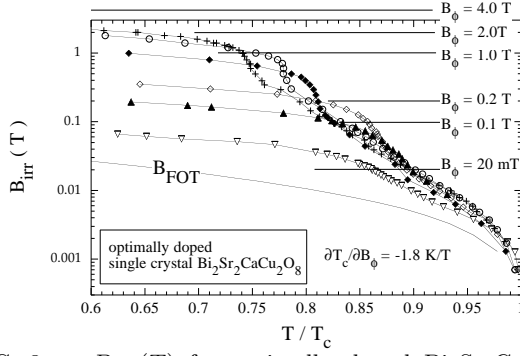


FIG. 3.  $B_{irr}(T)$  for optimally doped  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$  crystals with  $B_\phi$  between 0 and 4 T. The drawn line indicates the first order phase transition  $B_{FOT}$  in a pristine crystal.

again adopted,  $B_{irr} \propto (1 - T/T_c)^\alpha$  with  $\alpha \gtrsim 1$  (regime III). The behavior of the irreversibility line for different oxygen content is shown in Fig. 4. The exponential decrease at high  $T$  is steeper for the overdoped crystal, which has the smaller  $\lambda_{ab}$  and  $\gamma$ , and therefore the larger condensation energy  $\varepsilon_0/4\pi\xi^2$ , and the stronger intervortex- and vortex-defect interaction (both are proportional to the typical vortex energy scale  $\varepsilon_0 = \Phi_0^2/4\pi\mu_0\lambda_{ab}^2$ ;  $\xi$  is the coherence length).

In order to describe the low-field exponential temperature dependence of  $B_{irr}(T)$ , we exploit the fact that in this regime (I) vortex interactions are irrelevant to the column occupation [21]. In other words, each vortex line can become localized on an appropriate defect site. If a sufficient number of columns is available to every line (*i.e.* at large  $B_\phi$ ), extra free energy can be gained by the redistribution of pancakes constituting a given line over different columns. At low  $B_\phi \lesssim 0.5$  T this entropy gain is insufficient to balance the loss in vortex interaction energy, and pancakes belonging to the same line remain aligned on the same columnar defect. In either case, all pancakes are localized on a column. The superposition of the exponential portions of  $B_{irr}(T)$  for *all* matching fields  $30 \text{ mT} < B_\phi < 4 \text{ T}$  implies that at  $B_{irr}$  the redistribution over different column sites is irrelevant for the pancake delocalization mechanism — *i.e.* at and below  $B_{irr}(T)$ , pancakes belonging to the same vortex line are well aligned on the same site, belonging to a set of allowed “columnar-defect” sites. Then, we no longer need to consider other positions in the “intercolumn space” in our model description, which becomes that of a “discrete” superconductor. Since the allowed sites are more or less equivalent, the circumstance that they are in fact columnar defect sites becomes immaterial: namely, the free energy of all vortices is lowered by approximately the same amount. The low-field vortex state therefore does not differ fundamentally from that in the *unirradiated* crystal. Only, the vortex confinement in the defect potential, which plays the role of the “substrate potential” of Ref. [26], inhibits thermal line wandering and the elastic relaxation of the vortex lattice.

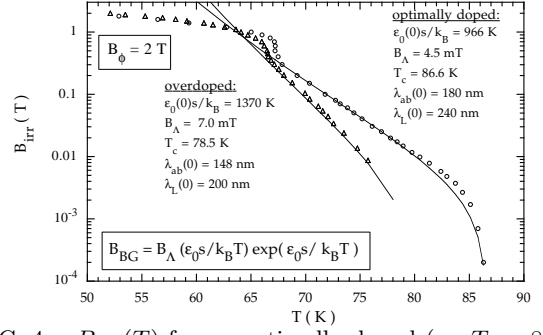


FIG. 4.  $B_{irr}(T)$  for an optimally doped ( $\circ$ ,  $T_c = 86.3$  K after irradiation) and a lightly overdoped ( $\triangle$ ,  $T_c = 78.5$  K)  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$  crystal, both with  $B_\phi = 2$  T. Drawn lines indicate fits to Eq. (2), with parameter values as indicated.

The main thermal excitations in this situation are expected to be small defects in the vortex lattice. In a Josephson-coupled layered superconductor, these amount to bound pancake vacancy-interstitial pairs within the same layer, *i.e.* the “quartets” of Ref. [27]. In the HII layered superconductor, such a pair corresponds to the “exchange” of one or more pancakes between two sites. The energy of the quartets is [27]

$$\varepsilon_q \approx 4c_{66}a_0^2s(R/\Lambda)^2 \approx \varepsilon_0s(R/\Lambda)^2, \quad (R \ll \Lambda) \quad (1)$$

with  $c_{66}$  the vortex lattice shear modulus,  $a_0 = (\Phi_0/B)^{1/2}$  the vortex spacing,  $R$  the distance between a bound vacancy and interstitial,  $\Lambda = [\lambda_{ab}^{-1} + (\gamma s)^{-1}]^{-1}$  the generalized penetration depth taking into account both magnetic and Josephson coupling, and  $\gamma s$  the Josephson length [18,27]. It was shown in Ref. [27] that the glass transition in a layered superconductor corresponds to the pair-unbinding transition. Correspondingly, the Bose glass transition is the unbinding temperature of the dislocation pairs in the “discrete” superconductor, *i.e.* the temperature above which pancakes can *diffuse* from line to line; it can be estimated as  $k_B T_{BG} = \varepsilon_q(R_l)$ , where  $R_l \approx n_l^{-1/2}$  and  $n_l$  is the equilibrium density of free dislocation pairs in the vortex liquid. Taking into account that only small pairs of size  $a_0$  (*i.e.* vacancies / interstitials) matter, one has  $n \approx a_0^{-2} \exp(-\varepsilon_0s/k_B T)$ ; the activation energy  $\varepsilon_0s$  is larger than that in the unirradiated superconductor because of the lack of lattice relaxation around the pair. Gathering terms,  $k_B T_{BG} = \varepsilon_0s(a_0/\Lambda)^2 \exp(\varepsilon_0s/k_B T_{BG})$  or

$$B_{BG} = B_\Lambda \left( \frac{\varepsilon_0s}{k_B T} \right) \exp \left( \frac{\varepsilon_0s}{k_B T} \right), \quad (B_\Lambda \ll B \ll B_\phi) \quad (2)$$

with  $B_\Lambda = \Phi_0/\Lambda^2$ ,  $\varepsilon_0$  and  $\Lambda$  to be evaluated at  $T_{BG}$ . The Bose-glass transition line does not depend on the details of the columnar defect potential such as pinning energy, column radius, or matching field. In Fig. 4, we compare Eq. (2) to the experimental results for  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$  with different oxygen content. We obtain excellent *quantitative* agreement using  $\lambda_{ab}$ -values from the literature

[22] and the single free parameter  $\gamma$  (contained only in  $B_\Lambda$ ). The values  $\gamma = 360$  and  $550$  found for the overdoped and the optimally doped material respectively are well within accepted experimental limits. Thus, Eq. (2) reproduces the correct field, temperature,  $B_\phi$ , and doping-dependence of the Bose-glass delocalization line.

The Bose-glass transition separates the low- $T$  phase in which pancake vortices can wander between columnar defects, but always remain bound to the same site, or vortex line, from the high  $T$ -phase in which diffusion of pancakes between sites (lines) is possible. Below  $T_{BG}$ , individual pancakes (“defect pairs”) cannot provide flux transport at low currents; vortex lines can only move as a whole, leading to the line-like behavior observed in Ref. [3–7]. Oppositely, the pancake diffusion above  $T_{BG}$  not only implies an Ohmic resistivity [5,6], but also that the linear nature of the vortex lines should no longer be apparent in the angular dependence of the transport properties. The intersite diffusion of pancakes (or “pancake-exchange”) means that at  $T > T_{BG}$  vortex lines are effectively entangled, on a scale  $l = s \exp(\varepsilon_0 s / k_B T)$ . The upper limit of the Bose-glass transition in layered superconductors is thus analogous to that in moderately anisotropic compounds in that it represents the boundary at which vortices become delocalized into an entangled flux liquid and where the angular dependence of the columnar defects’ contribution to the resistivity vanishes.

The meaning of  $B_\Lambda$  becomes apparent from Eq. (1), where  $\Lambda$  appears as the typical interaction distance between small dislocation pairs. For separations  $\gg \Lambda$ , *i.e.* for matching fields  $B_\phi \ll B_\Lambda$ , the formation of quartets demands a (shear) energy cost that is much greater than either  $\varepsilon_0 s$  or  $k_B T$ . The excitation of a pancake onto another site is then unlikely even in the vortex liquid. Delocalization will occur as it does in a moderately anisotropic superconductor: thermal wandering of the vortices into the intercolumn space lowers their binding energy until they can be liberated from the columns by thermal activation [4]. Similarly, in regime II ( $B > \frac{1}{6} B_\phi$ ), where the number of available tracks is insufficient not because of the intercolumn distance but because of track occupation by other vortices, delocalization is initiated by pancake wandering into the intercolumn space. Then,  $B_{irr}(T)$  increases with decreasing  $B_\phi$  because the number of available sites decreases. The *experimentally* found condition  $B < \frac{1}{6} B_\phi$  on the validity of Eq. (2) also describes its validity range at low  $B_\phi$ :  $B_{irr}$  becomes ion-dose dependent when  $B_\phi < B_\phi^{min} = 6B_\Lambda$ , in excellent agreement with the experimental values  $B_\phi^{min} \sim 30$  mT and  $B_\Lambda = 7.0$  mT (4.5 mT) found for overdoped (optimally doped)  $\text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8$  respectively.

Summarizing, in the regime of individual vortex line pinning by columnar defects in a layered superconductor, the upper limit of the Bose-glass transition corresponds to the onset of pancake vortex diffusion between different

lines and flux-line entanglement. However, in both the glass and the vortex liquid phases, *only* columnar defect sites are available to the pancakes; hence, a modelisation in terms of a “discrete superconductor” is appropriate. The onset of diffusion in this “discrete superconductor” presents a realization of the glass transition by unbinding of small defect pairs proposed in Ref. [27]; in the original continuous problem of an unirradiated layered superconductor, the defect-unbinding mechanism is masked by strong Gaussian thermal fluctuations [18].

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